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Direct generation of a Majorana mass for the neutron from exotic instantons

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ABSTRACT

We discuss a new mechanism in which non-perturbative quantum gravity effects directly generate a Majorana mass for the neutron. In particular, in string theory, exotic instantons can generate an effective six quark operator by calculable mixed disk amplitudes. In a low string scale scenario, with $M_S \simeq 10 \div 10^5$ TeV, a neutron–antineutron oscillation can be reached in the next generation of experiments. We argue that protons and neutralinos are not destabilized and that dangerous FCNCs are not generated. We show an example of quiver theories, locally free by tadpoles and anomalies, reproducing MSSM plus a Majorana neutron and a Majorana neutrino. These models naturally provide a viable baryogenesis mechanism by resonant RH neutrino decays, as well as a stable WIMP-like dark matter.

1. Introduction

Testing low energy B/L-violating processes is crucially important for our understanding of particle masses and matter–antimatter asymmetry in our Universe. Inspired by these deep motivations, experiments on neutrinoless-double-beta-decays are very active, constraining the Majorana neutrino. However, also a *neutron* can also have an effective Majorana mass term! Majorana himself first suggested the neutron as a Majorana fermion [1]. In terms of Weinberg effective operators, such a mass term corresponds to a six-quark dimension-9 term $(udd)^2/\mathcal{M}_{n\bar{n}}^5$. $\mathcal{M}_{n\bar{n}}, \tau_{n\bar{n}}, \delta m_{n\bar{n}}$ are connected to each other

$$\delta m_{n\bar{n}} = \tau_{n\bar{n}}^{-1} \simeq \left(\frac{\Lambda_{QCD}^6}{\mathcal{M}_{n\bar{n}}^5}\right) \simeq 10^{-25} \left(\frac{1000\ \text{TeV}}{\mathcal{M}_{n\bar{n}}}\right)^5 \text{eV} \tag{1}$$

where $\delta m_{n\bar{n}}$ is the Majorana mass of the neutron, $\tau_{n\bar{n}}$ the neutron–antineutron transition time.[1,2] Contrary to neutrini, a Majorana neutron can be directly tested in oscillations: neutron–antineutron transitions! The best limit on neutron–antineutron transitions in vacuum is around $\tau_{n\bar{n}} > 3$ yr, from Baldo-Ceolin experiment in Grenoble ('97') [6]. This seems surprising if compared to other rare processes like proton decay $\tau_p > 10^{35}$ yr and neutrinoless double-beta decays $\tau_{0\nu\beta\beta} > 10^{25}$ yr.[3] As discussed in [9–11], there is the exciting opportunity to enhance current limits on $n - \bar{n}$ transitions up to $\tau_{n\bar{n}} \simeq 300$ yr, with an experimental set-up *a la* Baldo-Ceolin – with external magnetic field $|\boldsymbol{B}| \sim 10^{-5} \div 10^{-6}$ Gauss.[4]

Recently, we have suggested that a Majorana mass for the neutron can be indirectly generated by non-perturbative effects of string theories known as exotic instantons [15–21]. Exotic instantons are peculiarly different from gauge instantons. In fact, they cannot be reconstructed by an ADHM classification of gauge instantons. Usually gauge instantons can 'strongly' violate axial

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[1] The validity of these estimations of non-perturbative QCD effects was checked in [2–5].

[2] Neutron–antineutron transition can be also a hint for a new fifth force interaction [7].

[3] A so fast $n - \bar{n}$ transition in vacuum does not destabilize nuclei: Neutron–antineutron oscillations are not just excluded by Superkamiokande experiment [8]. In fact, contrary to decay processes, a $n - \bar{n}$ oscillation is strongly suppressed by the nuclear binding energy in the nuclear environment. In fact, a transition from a binding neutron to a practically unbounded antineutron inside nuclei is energetically unfavored: the effective low energy Hamiltonian of the neutron–antineutron system will have diagonal terms with a difference between of 10^{33} digits up than off-diagonal Majorana masses, so that the transition time in nuclei will be suppressed up to $\tau_{n\bar{n}}^{Nuclei} > \tau_{exp}^{Nuclei} \simeq 10^{32}$ yr [8–11]. See Appendix C for explicit calculations of these suppression effects.

[4] To realize these experiments with suppressed magnetic fields is necessary in order to not suppress $n - \bar{n}$ oscillations. For example, the Earth magnetic field (0.5 Gauss) will split the energies of neutron and antineutron of $2\mu_n B_E \simeq 10^{-11}$ eV. This energy is 10^{12} higher than the present limit on the Majorana mass for the neutron ($\delta m_{n\bar{n}} < 10^{-23}$ eV $\simeq 10^{-8}$ s^{-1}).

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global symmetries. On the other hand, exotic instantons can 'strongly' violate global vector-like symmetries, like Baryon/Lepton ones. These effects are often calculable and controllable in open string theories. In open-string theories, all instantons admit a simple "geometric" interpretation. In fact they are Euclidean D-branes, or E-branes, wrapping an internal cycle, that could intersect the 'physical' D-branes. All gauge instantons of ADHM can be re-obtained by E-branes wrapping the same n-cycles of ordinary D-branes. On the other hand, 'exotic instantons' are E-branes wrapping *different* n-cycles of ordinary D-branes. See [12–14] for useful references on these aspects. As shown in [15–21], they can dynamically violate starting discrete symmetries Z_k of a perturbative Lagrangian. However, depending on their intersections with ordinary D-branes, they induce only specific operators, not necessarily all the possible Z_k-violating ones. In our first models, we have suggested that a $\mathcal{O}_{n\bar{n}}$ can be mediated by exotic matter with non-perturbative couplings induced by exotic instantons. In this sense, we have defined these mechanisms 'indirect' ones.

In this paper, we suggest a new mechanism for the generation of an effective Majorana mass for the neutron: we propose a simple and calculable mechanism generating $\mathcal{O}_{n\bar{n}}$ *totally* by exotic instantons, without the need for any mediator-fields. This could be a counter-example to UV Wilsonian completion of effective operators. Wilsonian UV completion has manifested itself as a successful approach in a lot of well understood examples in particle physics. Probably the most famous example is the UV completion of the Fermi model of weak interaction with a gauge theory of electroweak interactions, *i.e.* the non-renormalizable four fermion interaction is resolved as an exchange of a massive W boson. However, we will show how effective operators as six-quark operators can be completed by exotic instantons rather than by integrating-out new massive fields. In our proposal, we will show how for $M_S = 10 \div 10^5$ TeV, a neutron–antineutron transition can be found in the next generation of experiments. We will see how this scenario can be compatible with two of the most elegant solutions to the hierarchy problem of the Higgs mass: TeV-scale SUSY, or alternatively with low scale string theories with $M_S = 10$ TeV [22–24]. We are in LHC era, which will provide a lot of important inputs also for our model in the immediate future! On the other hand, we will also consider 'less appealing' scenarios for LHC, compatible with a PeV-ish $\mathcal{M}_{n\bar{n}}$. In fact, the case of a $M_{SUSY} \simeq M_S \simeq 10^3$ TeV remains compatible with a $\mathcal{M}_{n\bar{n}} \simeq 1000$ TeV neutron–antineutron transition. In this last case, electroweak scale is fine-tuned for a factor $m_H/M_S \sim 10^{-4}$, rather than $m_H/M_{Pl} \sim 10^{-17}$, *i.e.* the hierarchy problem is alleviated. We will also discuss explicit examples of un-oriented quiver theories reproducing at low energy limit a (MS)SM with a Majorana neutron, a electroweak-scale mu-term, neutrini masses from Right-handed neutrini (compatible with RH-neutrino decayRal baryogenesis), WIMP-like dark matter, without destabilizing nucleons.

2. (MS)SM quivers for an exotic Majorana neutron

A MSSM can be embedded in a quiver theory,[5] with at least three nodes, reproducing a gauge theory $U(3)_a \times U(2)_b \times U(1)_c$, in type II A string theory. The basic elements are few ones: $D6$-branes wrapping 3-cycles in the Calabi–Yau CY_3, one of these will be a flavor brane, one Ω-plane, $E2$-instantons, open strings attached to $D6$-branes and $E2$-branes. Stacks of three parallel D6-branes will reproduce $U(3)_a$, and so on. (MS)SM matter fields in the bi-fundamental representations of SM gauge groups are reproduced, in the low energy limit, by open (un)oriented strings attached to two intersecting D6-branes' stacks. For example, a Q_L superfield comes from an open string attached to a-stack, reproducing $U(3)$, and b-stack, reproducing $U(2)$. The number of generations is reproduced by the number of intersections among D6-branes' stacks. For example, the three generations of quarks correspond to three intersections between a-stack and b-stack. In these models, hypercharge $U(1)_Y$ is a massless combination of $U(1)_a$ contained in $U(3)_a \simeq SU(3) \times U(1)_a$; $U(1)_b$ contained in $U(2)_b \simeq SU(2) \times U(1)_b$; and $U(1)_c$:

$$U(1)_Y = q_a U(1)_a + q_b U(1)_b + q_c U(1)_c \tag{2}$$

As regards the two extra anomalous $U(1)$'s, they can be cured by the generalized Green–Schwarz mechanism. Intriguingly, anomalous $U(1)$ are impossible to be consistently included in gauge theories while they can be opportunely cured in string theories. In the stringy extension of the (MS)SM as the ones in consideration, two new vector bosons Z', Z'' are generically predicted, getting a mass through a Stückelberg mechanism, and interacting through generalized Chern–Simon (GCS) terms. In fact GCS are introduced in order to cancel all anomalies. See [53–60] for an extensive discussion on these aspects.

In a minimal three-nodes configuration, (MS)SM superfields can be reconstructed as follows: Q_L as $(a,\bar{b})$ or (a,b); U_R as $(\bar{a},\bar{c})$; D_R as $(\bar{a},c)$; H_u as (b,c); H_d as $(\bar{b},c)$; L as $(\bar{b},\bar{c})$ or $(\bar{b},\bar{c})$; E_R as (c,c').

For these intersections, the correspondent hypercharge is $U(1)_Y = \frac{1}{6}U(1)_a + \frac{1}{2}U(1)_c$.

However, in order to guarantee that such a model is free by tadpoles and that $U(1)_Y$ associated to Y is a massless combination, usually one has to add extra exotic matter in order to satisfy these two stringent condition. A complete classification of extra massive states was shown in [44]. Typically, extra vector-like pairs and charged singlets are often introduced for consistency.

These models cannot reproduce all the MSSM Yukawa couplings at perturbative level. For example, $y_u H_u Q U^c$ is generate at perturbative level, but not $y_d H_d Q D^c$ and $y_l H_d L E^c$. However, they can be generated by $E2$-instantons, opportunely intersecting D6-branes' stacks, with $y_d \sim e^{-S_{E_2^d}}$ and $y_l \sim e^{-S_{E_2^l}}$ [61–63], where $S_{E_2^d, E_2^l}$ depend on geometric moduli associated to 3-cycles wrapped on CY_3 by E_2^d, E_2^l-instantons. Also a μ-term is not generated at perturbative level, but it can be generated by exotic instantons as just proposed in [61–63]. In particular, as shown in [61,62], one obtain $\mu = e^{-S_{E2''}} M_S$ where M_S is the string scale, while $e^{-S_{E2''}}$ depends by geometric moduli parameterizing 3-cycles wrapped on CY_3 by $E2''$-branes respectively. Practically, in these local models, we can consider $e^{-S_{E2'',E2^{d,l}}}$ as free-parameters, depending on the particular geometry of the exotic instantons considered.

Let us consider in this class of models, the presence of a new $E2$-brane intersecting two times the $U(3)$-stack, two times the $U(1)$-stack, four time the $U(1)'$-stack. These mixed disk amplitudes lead to effective interactions between U^c, D^c and fermionic zero moduli (modulini). In fact, modulini are obtained by open strings attached to $D6$-stacks and $E2$-stacks rather than to $D6$–$D6$. Let us assume that this $E2$-instanton has a Chan–Paton factor $O(1)$, *i.e.* it is placed on a Ω^+-plane (symmetric). Calling τ^i modulini living between $U(3)$–$E2$, α modulini between $U(1)$–$E2$ and β between $U(1)'$–$E2$, the following effective interactions are generated:

$$\mathcal{L}_{eff} \sim c_f^{(1)} U_f^i \tau_i \alpha + c_f^{(2)} D_f^i \tau_i \beta \tag{3}$$

Integrating on the modulini space associated to the D6–E2 intersections, we obtain

[5] See for useful references on open string theories and orientifolds [25–35]. See [36–52,64–70] for useful literature on MSSM quiver theories.

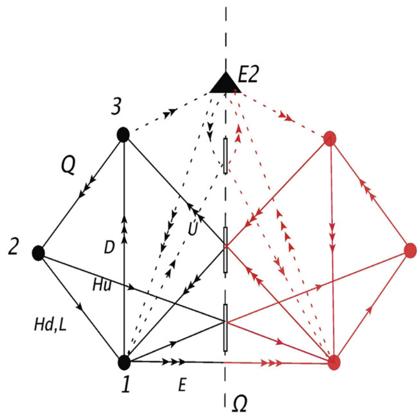

Fig. 1. A three-nodes quiver for a Majorana neutron. We do not report all instantons generating Yukawa couplings $y_l H_d L E^c$ and $y_d H_d Q D^c$, and the μ-term. We also omit possible exotic fields getting consistent our quiver from the point of view of tadpoles cancellations and massless hypercharge $U(1)_Y$, discussed in [44]. This quiver shows the number of intersections of $E2$ with 3,2,1-stacks, generating six-quarks superpotentials.

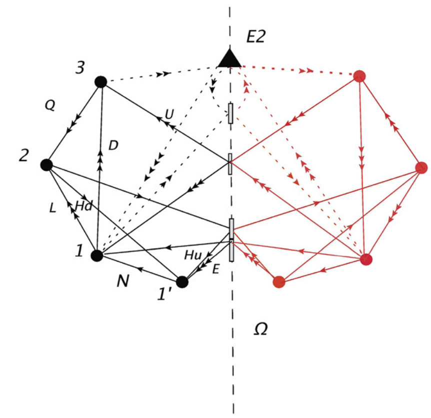

Fig. 2. A four-nodes Madrid quiver for a Majorana neutron. We do not report all instantons generating Yukawa couplings, μ-term and RH neutrini masses. We also omit possible exotic fields getting consistent our quiver from the point of view of tadpoles cancellations and massless hypercharge $U(1)_Y$, discussed in [44]. This quiver shows the number of intersections of $E2$ with 3,2,1-stacks, generating six-quarks superpotentials.

$$\mathcal{W}_{E2} = \int d^6\tau d^4\beta d^2\alpha e^{\mathcal{L}_{eff}}$$
$$= \mathcal{Y}^{(1)} \frac{e^{-S_{E2}}}{M_S^3} \epsilon_{ijk}\epsilon_{i'j'k'} U^i D^j D^k U^{i'} D^{j'} D^{k'} \tag{4}$$

where $\mathcal{Y}^{(1)}_{f_1 f_2 f_3 f_4 f_5 f_5} = c^{(1)}_{f_1} c^{(1)}_{f_2} c^{(2)}_{f_3} c^{(2)}_{f_4} c^{(2)}_{f_5} c^{(2)}_{f_6}$ is the flavor matrix, and $c^{(1,2)}$ are related on the particular homology and topology of the mixed disk amplitudes, so that we can assume these as free parameters. Superpotential (4) corresponds to the ope for $n - \bar{n}$ transitions

$$\mathcal{O}_{n\bar{n}} = \frac{y_1}{\mathcal{M}^3_{E2} M^2_{SUSY}} (u^c d^c d^d)(u^c d^c d^c) \tag{5}$$

where $\mathcal{M}^3_{E2} = e^{+S_{E2}} M_S^3$, and M_{SUSY} comes from the SUSY conversion of squarks into quarks through the exchange of a gaugino as a gluino, zino or photino, $y_1 = \mathcal{Y}^{(1)}_{111111}$.

For the quiver shown in Fig. 1, a first problem is the perturbative generation of $\mathcal{W}_{\Delta L=1} = y_{LQU} L Q U^c$, violating L of $\Delta L = 1$. This superpotential has to be tuned close to zero, in order to avoid a dangerous proton destabilization.

Now, let us discuss another possible case with four nodes, known as Madrid-embedding, with hypercharge $U(1)_Y = \frac{1}{6}U(1)_a + \frac{1}{2}U(1)_c + \frac{1}{2}U(1)_d$. In this class of quiver, one could obtain discrete symmetries like R-parity from the Stückelberg mechanism of anomalous $U(1)$s, as discussed in [71]. For this motivation, the generation of a Majorana mass for the neutron, as well as a Majorana mass for the neutrino, comes from an opportune $E2$-instantons, dynamically breaking R-parity so that a selection rule $\Delta B = \Delta L = 2$ will emerge. As a consequence, no other bilinear or trilinear R-parity violating superpotentials are generated.

A Madrid-embedding allows for Q_L as $(a, \bar{b})$ or (a, b); U_R as $(\bar{a}, \bar{c})$ or $(\bar{a}, \bar{d})$; D_R as $(\bar{a}, c)$ or $(\bar{a}, d)$; L_L as $(b, \bar{c})$, $(\bar{b}, \bar{c})$, $(b, \bar{d})$, $(\bar{b}, \bar{d})$; E_R as (c, d) or A_c or S_c; N_R as A_b or $\bar{A}_b$ or $(c, \bar{d})$ or $(\bar{b}, d)$; H_u as $(\bar{b}, c)$ or (b, c) or $(\bar{b}, d)$ or (b, d); H_d as $(b, \bar{c})$ or $(\bar{b}, \bar{c})$ or $(b, \bar{d})$ or $(\bar{b}, \bar{d})$.

Generically, also in this case, several MSSM Yukawa couplings will be not generated at perturbative level, but they can be non-perturbatively generated by opportune $E2$-instantons, as mentioned above for three-nodes' quivers.

In Fig. 2, we show a possible example, in which MSSM Yukawa couplings have to be generated by exotic instantons. In order to cancel all tadpoles and to guarantee a massless hypercharge, one has to introduce extra exotic matter, as in the case mentioned above. A complete classification of exotic superfields introduced for consistency was shown in [44] also for this case. In Fig. 2, we show how an $E2$-instanton generating (4) can be easily introduced as in the three node case.

These examples seem to sustain the quite generality of such a mechanism, for several models in literature [39,52].

In other models a possible viable alternative can be considered: an $E2$-instanton, with modulini coupled to Q and D rather than U and D, as (4)

$$\mathcal{L}_{eff} \sim c^{(3)}_f Q^i_{f,\alpha} \gamma_i \eta^\alpha + c^{(4)}_f D^i_f \gamma_i \zeta \tag{6}$$

where γ_i are modulini living between $E2$–$U(3)$, η between $E2$–$U(2)$, ζ between $E2$–$U^I(1)$ ($U(1)^I$ is the involved $U(1)$ depending by the specific quiver). Integrating on the modulini space we generate an effective superpotential

$$\mathcal{W}_{E2} = \int d^6\tau d^2\zeta d^8\eta e^{\mathcal{L}_{eff}}$$
$$= \mathcal{Y}^{(2)} \frac{e^{-S_{E2}}}{M_S^3} \epsilon_{ijk}\epsilon_{i'j'k'} \epsilon^{\alpha\alpha'} \epsilon^{\beta\beta'} Q^i_\alpha Q^j_\beta D^k Q^{i'}_{\alpha'} Q^{j'}_{\beta'} D^{k'} \tag{7}$$

The associates ope for $n - \bar{n}$ transitions is

$$\mathcal{O}_{n\bar{n}} = \frac{y_2}{\mathcal{M}^3_{E2} M^2_{SUSY}} (q^c q^c d^d)(q^c q^c d^c) \tag{8}$$

where again $\mathcal{M}^3_{E2} = e^{+S_{E2}} M_S^3$ and $y_2 = \mathcal{Y}^{(2)}_{111111}$.

Finally, similarly to Figs. 1–2, the mechanism can be implemented in a completely consistent quiver without extra exotic colored or electroweak states (see Fig. 3) and with all SM Yukawa couplings perturbatively allowed,

The hypercharge in this model is the combination of 3 charges, coming from $U(1)_3$, $U(1)$ and $U'(1)$:

$$Y(Q) = c_3 q_3 + c_1 q_1 + c_1' q_1' \tag{9}$$

We can fix the coefficients in such a way as to recover the standard hypercharges:

$$Y(Q) = \frac{1}{3} = c_3 \tag{10}$$

$$Y(U^c) = -\frac{4}{3} = -c_3 - c_1' \tag{11}$$

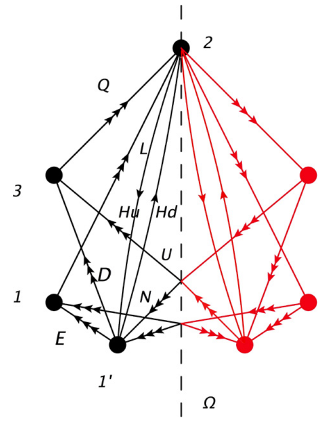

Fig. 3. A four-nodes SM quiver. We omit the E2-instanton generating a Majorana mass for the neutron. Similarly to Figs. 1–2 such an $E2$-instanton can be consistently introduced.

$$Y(D^c)=\frac{2}{3}=-c_3+c_1' \tag{12}$$

$$Y(L)=-1=c_1 \tag{13}$$

$$Y(H_d)=-Y(H_u)=-1=-c_1' \tag{14}$$

$$Y(E^c)=2=-c_1+c_1' \tag{15}$$

$$Y(N_R)=0=-c_1-c_1' \tag{16}$$

leading to

$$c_3=\frac{1}{3},\;\; c_1=-1,\;\; c_1'=1 \quad\to\quad Y=\frac{1}{3}q_3-q_1+q_{1'} \tag{17}$$

Let us note that the quiver in Fig. 3 is free of tadpoles and the hypercharge $U(1)_Y$ is massless. A generic quivers has to satisfy two conditions in order to be anomalies'/tadpoles' free and in order to have a massless hypercharge. The first one associated to tadpoles' cancellations is

$$\sum_a N_a(\pi_a+\pi_{a'})=4\pi_\Omega \tag{18}$$

where $a=3,1,1'$ in the present case, π_a 3-cycles wrapped by "ordinary" D6-branes and $\pi_{a'}$ 3-cycles wrapped by the "image" D6-branes. Eq. (18) can be rewritten as a function of field representations

$$\forall\, a\neq a' \qquad \#F_a-\#\bar{F}_a+(N_a+4)(\#S_a-\#\bar{S}_a) + (N_a-4)(\#A_a-\#\bar{A}_a)=0 \tag{19}$$

where $F_a, \bar{F}_a, S_a, \bar{S}_a, A_a, \bar{A}_a$ are fundamental, symmetric, antisymmetric of $U(N_a)$ and their conjugate. For $N_a>1$ these coincide with the absence of irreducible $SU(N_a)^3$ triangle anomalies. For $N_a=1$, these are stringy conditions that can be rephrased as absence of 'irreducible' $U(1)^3$ *i.e.* those arising from inserting all the vector bosons of the same $U(1)$ on the same boundary. Let us explicitly check tadpole cancellation for the $3,1,1'$ nodes:

$$U(3):\quad 2n_Q-n_D-n_U=6-3-3=0 \tag{20}$$

$$U(1):\quad 2n_L-n_E-n_N=6-3-3=0 \tag{21}$$

$$U(1)':\quad n_E-n_N+3n_D-3n_U=3-3+3\cdot 3-3\cdot 3-3=0 \tag{22}$$

The quiver in Fig. 3 also satisfies the condition a massless vector boson $U(1)_Y$, with $Y=\sum_a c_a Q_a$,

$$\sum_a c_a N_a(\pi_a-\pi_a')=0 \tag{23}$$

corresponding to

$$c_a N_a(\#S_a-\#\bar{S}_a+\#A_a-\#\bar{A}_a) - \sum_{b\neq a} c_b N_b(\#(F_a,\bar{F}_b)-\#(F_a,F_b))=0 \tag{24}$$

For $c_3^Y=1/3$, $c_1^Y=-1$, $c_1'^{\,Y}=1$, these conditions are satisfied in each quiver nodes. On the other hand, the massive (anomalous) $U(1)$'s correspond to $3Q_3+Q_1$ and to $3Q_3-Q_1'$. This can be demonstrated calculating the anomaly polynomial.

2.1. Comments on proton decay, dark matter, FCNCs, baryogenesis

In our models, exotic instantons generate $\Delta L=2$ violating mass terms for neutrini and $\Delta B=2$ violating mass term for the neutron. As a consequence, in R-parity preserving models, R-parity is dynamically violated at non-perturbative level, so that other $\Delta L, B=1$ superpotentials are not generated. As a consequence, a selection rule $\Delta B, L=2$ has emerged in this mechanism. No-proton decays operators have been generated. In fact,

$$\mathcal{W}_{MSSM}+\frac{1}{2}m_N NN+\frac{1}{\mathcal{M}_{n\bar{n}}^3}U^cD^cD^cU^cD^cD^c \tag{25}$$

cannot generate a proton decay as well as the lightest neutralino remains stable. Clearly the same is valid for

$$\mathcal{W}_{MSSM}+\frac{1}{2}m_N NN+\frac{1}{\mathcal{M}_{n\bar{n}}^3}QQD^cQQD^c \tag{26}$$

On the other hand, extra diagrams CP-violating FCNCs in mesons' physics, such as $K_0-\bar{K}_0, B_0-\bar{B}_0,\dots$, are strongly suppressed in our model, up to 1000 TeV-scale. This is an interesting difference with respect to other neutron–antineutron models, usually involving extra colored states leading to non-negligible contributions to FCNCs.

The next answer regards baryogenesis. In fact, $(u^cd^cd^c)^2/\mathcal{M}_{n\bar{n}}^5$ or $(qqd^c)^2/\mathcal{M}_{n\bar{n}}^5$ can generate collisions $ud^cd^c\to\bar{u}^c\bar{d}^c\bar{d}^c$ or $qqd^c\to\bar{q}\bar{q}\bar{d}^c$. Supposing an initial ΔL generated by RH-neutrini decays one has to be careful about washing-out action of six-quarks' collisions $B-L$ and sphalerons $B+L$. However, the problem is solved if simply $m_N<\mathcal{M}_{n\bar{n}}$: in this way, six quarks collisions are not relevant during out-of-equilibrium RH-neutrini decays.

Taking into account these considerations, let us discuss the different branches of parameters in our model. In order to obtain a $\mathcal{M}_{n\bar{n}}\simeq 1000$ TeV, we have several region of parameters. We discuss the main interesting ones.

1) $M_{SUSY}\simeq\mathcal{M}_{E2}\simeq 1000$ TeV. $\mathcal{M}_{E2}=e^{+S_{E2}/3}M_S$, so that a case $M_{SUSY}\simeq M_S\simeq 1000$ TeV, with $e^{+S_{E2}/3}\simeq 1$,[6] seems a natural possibility. The condition $e^{+S_{E2}/3}$ is geometrically understood as small radii 3-cycles wrapped by the $E2$-brane in the CY_3. In this scenario, $M_N\simeq e^{-S_{E2'}}\times(1000\ \text{TeV})$. However, usually, for a successful baryogenesis, M_N mass has to be higher than Davidson–Ibarra bound $M_N>10^6$ TeV. On the other hand, RH neutrini masses are generated by $E2'$ and their values depend on the particular geometry of the mixed disk amplitudes. As a consequence, a *resonant leptogenesis* scenario seems favored by this space of the parameters. considering at least two highly degenerate RH neutrini masses. In this case, Davidson–Ibarra bound is completely avoided, and RH neutrini masses can be also $M_R\simeq 1$ TeV or so, as quantitatively shown in [72]. A RH neutrini degenerate spectrum can be

[6] However, in this limit, 3-cycles are so small that semiclassical approximation cannot be applied. Limit on semiclassical validity is e^{-1/g_s}. On the other hand, g_s can also be large as $g_s\simeq 0.5$ or so.

understood geometrically by mixed disk amplitudes involved. As a consequence, $e^{-S_{E2'}}$ is constrained to $1 \div 10^{-3}$. Hierarchy problem of the Higgs mass is alleviated of a 10^{-28} factor in this scenario, as $(M_S/10^{19}\ \text{GeV})^2$. In this scenario, LHC will not observe any signatures.

2) $M_{SUSY} \simeq 1$ TeV and $\mathcal{M}_E \simeq 10^5$ TeV. In this case, a scenario for supersymmetry at LHC can be immediately tested in the next run. Such a scenario corresponds to several different String-scales: $\mathcal{M}_E = M_S e^{+S_{E_2}/3} \simeq 10^5$ TeV is compatible with $M_S = 10$ TeV and $e^{+S_{E2}/3} \simeq 10^4$ (large 3-cycles), as well as with $M_S = 10^5$ TeV and $e^{+S_{E2}/3} \simeq 1$ (small 3-cycles). As regard leptogenesis, $M_N = e^{-S_{E2'}} M_S$ so that $M_N = 1 \div 10^3$ TeV for several different combination of $e^{-S_{E2'}}$ and M_S (*i.e.* $e^{+E2/3}$). Again, in all these scenarios, a resonant leptogenesis with at least two degenerate species of RH neutrini is desired. In this case a stable neutralino dark matter with mass $10 \div 1000$ GeV is easily obtained, and the mu-problem is solved and understood as a hierarchy generated by the $E2''$-brane among the string scale and μ as $\mu = e^{-S_{E2''}} M_S$.[7]

The most optimistic scenario for LHC is for $M_S = 10$ TeV. In this case, signatures of Stringy Regge states will be immediately tested by LHC. For $M_S = 10$ TeV, also signatures by anomalous Z', Z'' with Stückelberg masses $m_{Z',Z''} \sim M_S$ can be tested at LHC, with peculiar channels from Generalized-Chern–Simons terms. See [74] for a recent discussion on these aspects.

3. Conclusions and remarks

In this paper, we have discussed a simple mechanism directly generating a Majorana mass for the neutron from exotic instantons. These effects are completely calculable and controllable. Usually, a Weinberg operator is UV completed by massive ordinary fields, integrated-out at low energy limit. In our case, we have shown a counter-example to the Wilsonian UV completion: a six quark operator like $(udd)^2/M_{n\bar{n}}^5$ is UV completed by exotic instantons.[8]

We have explicitly discussed examples of quivers reproducing the (MS)SM, with a Majorana neutron and Majorana neutrini. In these models, operators leading to a proton decay are not generated. In this framework, dark matter is not destabilized, and a successful resonant leptogenesis through RH neutrini decays can be realized. Starting from a B/L preserving model, exotic instantons have dynamically broken R-parity, so that $\Delta(B-L) = 2$ selection rules have naturally emerged, rather than imposed *ad hoc*. In the 'LHC era', LHC data will provide important inputs for our model, constraining or likely discovering direct signatures also connected to our class of models, for several different regions of the parameters' space... On the other hand, the possibility to improve the best current limit on neutron–antineutron transition in the near future is technically possible and well motivated by theoretical principles. The 'crazy idea' of Majorana in '37' has never been so relevant and intriguing as now!

7 Alternatively, the presence of a parallel intersecting D-branes' world can contribute to dark matter. If the vev in the parallel sector was different from the vev in our sector, the dark halo would be a non-collisional one, composed of dark atoms. See [73] for a recent discussion on these aspects and implications in dark matter direct detection phenomenology.

8 We are tempted to suggest that exotic instantons can be viewed as classicalons in internal dimension from the point of view of scattering amplitudes. Classicalization was firstly suggested by Dvali and collaborators [75,76], and recently it was studied considered in the contest of non-local QFTs [77] (see also [78] for discussions of scattering amplitudes in $\mathcal{N} = 1$ non-local QFTs). However, this conjecture will deserve future investigations beyond the purposes of these paper.

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